LABORATORY EXPERIMENTS ON MAGNETIC RECONNECTION AND CURRENT SYSTEMS


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ABSTRACT

After a brief review of laboratory experiments involving magnetic reconnection a series of basic physics experiments on reconnection phenomena is described. These include magnetic annihilation, transport of magnetic energy by waves, stable and unstable current sheets, energy conversion mechanisms, and the role of global current systems vs. local reconnection processes. The present laboratory experiment deals with current systems driven by electric fields resulting in particle flows. Currents are dominantly carried by electrons. The self-consistent current sheet thickness is usually smaller than an ion Larmor radius. Consequently, MHD concepts such as frozen-in magnetic fields, wave transport at the Alfvén speed, and absence of electrostatic electric fields break down. Instead, the dynamics of reconnection is controlled by the fast speed of electrons, i.e., magnetic fields propagate at whistler wave speeds, current sheets tear by electron tearing modes, ions are accelerated by space charge electric fields, and anisotropic particle distributions are common occurrences. The role of a magnetic field component B_y along the separator has been investigated. Although three-dimensional reconnection with kinetic effects is a difficult topic for theory and simulations it is likely to be important in various space and laboratory plasmas.

INTRODUCTION

Reconnection processes occur in many laboratory devices in a wide range of plasma parameters. In a possible classification one can distinguish those devices which are specifically built to study reconnection from others built for magnetic confinement of fusion-type plasmas. The former group includes linear plasma columns with external conductors (multipoles) /1-4/, pinch plasmas with magnetic null points /5-7/, and streaming plasmas interacting with magnetic dipoles modeling the earth’s magnetosphere (terrella) /8-10/. The latter group includes predominantly tokamaks /11-12/, spheromaks /13,14/ and reverse field pinches /15,16/. The motivation to study reconnection processes varies considerably in the different devices. Although the changes in magnetic field topologies is of common interest in all reconnection experiments the basic research is concentrated on the energy conversion processes while the fusion research is focused on confinement and transport processes associated with reconnection. Anomalous resistivity and its effects on reconnection rates has been studied in linear devices /17,18/ while magnetic helicity conservation is a major topic of toroidal reverse field pinches /19,20/.

Due to the various aspects of reconnection there is probably no single laboratory device in which all questions can be investigated at once. The large fusion devices with high magnetic Reynolds numbers are best suited to study MHD modes but since in-situ magnetic measurements are not feasible the fine-scale structure of magnetic null regions remains as elusive as in space observations. The smaller basic laboratory devices operating at lower density and temperature permit detailed in-situ probing of current sheets hence, are best suited to study non-MHD processes in the null regions. While in principle, such research devices could be scaled up to large magnetic Reynolds numbers the expense of generating a large volume of magnetized plasma is beyond the range of most basic research projects. Thus, the physics of reconnection has to be composed of observations from various laboratory devices. The present contribution comes from a basic linear plasma device and addresses various non-MHD processes which are not predicted by standard reconnection theories. The breakdown of MHD occurs when the ions become unmagnetized near magnetic null points which occurs over a much wider region than the classical diffusion region (R_m = 1). The consequence is a decoupling of the magnetic field from the mass of the ions to that of the light electrons. Magnetic energy is now transported by fast whistler waves rather than slow Alfvén waves, and the direction of energy flow can be predominantly along the separator provided there is a field component B_y. Space-charge electric fields build up, double layers can form at large electron drifts (v_d - v_e), and non-Maxwellian particle distributions can arise from acceleration in localized electric fields. Thus, the null region is a highly dynamic region in comparison with the simple
diffusion region extrapolated from SHE theory.

Another important aspect of reconnection learned in the laboratory is the importance to consider the current closure, i.e. all properties of the global current circuit. Except for the exceptional case of a uniform current path such as for a symmetric torus (e.g. tokamak), the reconnection electric field $E = -\partial A/\partial t$ is not uniformly distributed along the separator as assumed in 2D-theories. Observations have shown that the global reconnection voltage $V = -\partial \phi/\partial t$ drops off in regions of low conductivity (low density, constrictions, normal magnetic fields, boundaries, etc.) which may lie in the current closure path rather than in the current sheet itself. Hence, the rate of reconnection and energization may be controlled by plasma parameters outside of the null region. Magnetic energy is transported to the dissipation regions along the neutral line by fast whistler modes. A search for reconnection signatures (electric fields $E_n$, accelerated particles) in the neutral sheet then proves futile in spite of an ongoing reconnection process. Thus, in general, the events along the third dimension ($y$-axis in magnetospheric coordinates) are as important to consider as those in the $x$-$z$ plane.

The paper is organized in several sections, starting with basic experiments such as diffusion and wave transport of fields, followed by properties of neutral sheets, tearing modes, and acceleration processes. Only highlights of the experiments are given while details are found in the quoted references.

MAGNETIC DIFFUSION

Most classical papers on reconnection assume that in the immediate vicinity of a magnetic neutral sheet the process of magnetic diffusion is responsible for field line reconnection. Experimentally, the diffusion process can be isolated in an arrangement shown schematically in Figure 1. A large afterglow plasma column is generated by a pulsed dc discharge. An insulated thin wire is suspended along the axis parallel to the static magnetic field $B_0$. A current $I(t)$ is pulsed through the wire and the associated magnetic field $B(r,t)$ is measured with a small movable magnetic probe during repeated discharge pulses. Due to the simplicity of the coaxial arrangement the magnetic field has only one component which varies only in one direction and in time, $B_\phi(r,t)$. The penetration of the magnetic field is controlled by the vector diffusion equation $\nabla \times B = \sigma \partial B/\partial t$ which reduces to a scalar equation

$$\frac{\partial B}{\partial t} = \sigma \nabla \times B$$

from the measurement of $B_\phi(r,t)$ one can not only determine the induced plasma current density $J = \nabla \times B/\mu$ but also the diffusivity $D = \sigma \nabla \times B / \sigma$ or conductivity $\sigma$ by solving the vector diffusion equation. These quantities yield the Ohmic electric field $E = J/\sigma$ which is often difficult to measure directly in highly conducting plasmas.
Figure 2a shows the measured magnetic field $B_\Phi$ vs. radius $r$ at different times $t$ after penetration of magnetic fields into plasma.

The penetration of magnetic field $B_\Phi(r,t)$ into the plasma is governed by the vector diffusion equation. The slow penetration is caused by induced shielding currents in the plasma, a typical trace of which is shown in Figure 2b at a fixed time, $t = 5 \mu$s.

In time the cylindrical current layer diffuses away from the wire, broadens in width and decays in magnitude, coincident with the complete penetration of the magnetic field into plasma.

The reverse process, i.e. the decay of an embedded magnetic in a plasma occurs upon switch-off of the wire current and is shown in Fig. 3. Although the wire current is switched off at $t = 0$ the magnetic field in the plasma is maintained by induced plasma currents $J_2$ which are again distributed on expanding cylindrical shells as in the turn-on case, except that now the

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**Fig. 2.** (a) Radial magnetic field profiles at different times $t$ after switch-on of a current in a wire at $r = 0$. The penetration of magnetic field $B_\Phi(r,t)$ into the plasma is governed by the vector diffusion equation.
Fig. 2. (b) Typical current density profile \( J_z(r,t=S\mu\text{sec}) \). Insert shows schematically that the shielding current flows in cylindrical shell antiparallel to the wire current. In time, the current shell expands in radius and thickness and decreases in magnitude.

current direction is reversed. The time scale for the magnetic fields to decay depends on scale length and conductivity \( (t = \mu \sigma r^2) \). Since the primary circuit is open \( (I = 0) \) none of the stored magnetic field energy is transferred back to it, hence it is entirely dissipated or annihilated in the plasma. Small currents were used in order to prevent significant Joule heating of electrons which would locally vary the Spitzer conductivity and result in a complicated nonlinear diffusion process, \( \sigma(r,t,B) \). The diffusion process observed dominates over any wave transport, i.e. there is no radial or axial oscillatory \( B \)-field configuration excited.

MAGNETIC ENERGY TRANSPORT BY WAVES

Ideal MHD theory predicts that dc magnetic fields are “frozen in” a perfectly conducting plasma \( (B_0 = 0) \) and can change only by fluid motion. However, small amplitude magnetic perturbations can propagate as Alfvén waves and thereby transport magnetic energy. In the non MHD regime of neutral sheets the corresponding mode is the electron whistler wave. The following laboratory observations have shown that this mode controls the speed of current propagation and magnetic energy flow in the regime of unmagnetized ions.
Fig. 3. Radial magnetic field profiles at different times $t$ after switch-off of the wire current. The initial magnetic field decays by diffusion (magnetic field annihilation). The current density profiles $J_z(r,t)$ are as during switch-on but with reverse sign.

Fig. 4. Simplified experimental setup for generating a field-aligned current system between two electrodes in a magnetoplasma in order to measure the propagation speed of currents in plasmas. The current front propagates at the speed of a whistler wave packet irrespective of whether the electrodes draw ions, electrons, or emit a fast electron beam.
Figure 4 shows a simple experimental arrangement whereby current is drawn between two electrodes along a uniform magnetized afterglow plasma /21-23/. The circuit is completed via a battery, switch and external conductor. Measured are the magnetic field perturbations caused by current flow in the plasma. Complete space-time dependence of $\mathbf{B} = \mathbf{A} \times \mathbf{E}$ yields the current density $\mathbf{J} = \nabla \times \mathbf{A} / \nu$. The observations focus on the speed of the current front ($v_1$) and the field structure at the front. The results show that the perturbation is carried by whistler waves irrespective of whether the current is established by ion or electron collection or fast electron emission at the electrodes.

CURRENT DENSITY PROPAGATION/GROWTH

Fig. 5. Contours of constant current density in x-z planes at different times $t$ after turn-on of the potential to a 2 cm diam. disc electrode at $z = 0$ drawing electrons along $\mathbf{E}_0$. Dashed line connects the fronts of the same low current contours and defines the propagation speed $v_{11} = \Delta z / \Delta t$. 

Fig. 5. Contours of constant current density in x-z planes at different times $t$ after turn-on of the potential to a 2 cm diam. disc electrode at $z = 0$ drawing electrons along $\mathbf{E}_0$. Dashed line connects the fronts of the same low current contours and defines the propagation speed $v_{11} = \Delta z / \Delta t$. 
Figure 5 depicts the current density $J(x,z)$ as contour plots in planes along (z) and across (x) the dc magnetic field $B_0 (= 300 \text{ G})$ at different times $t$ after switch-on of the voltage $V (= 100 \text{ V})$ to a circular ($r = 1 \text{ cm}$) plane cold electrode centered at $x=y=z=0$. The front of a typical contour of constant current density penetrates into the plasma at a speed of $\Delta z/\Delta t = 6 \times 10^5 \text{ cm/sec}$ as indicated by the dashed line. The speed increases with magnetic field ($\propto B^{1/2}$) and decreases with density ($\propto n^{-1/2}$), is independent of charged particle species, temperature or current density, hence corresponds to a wave speed rather than a particle drift speed.

The dispersive property of the current/magnetic field front is shown in Fig. 6 which displays the magnetic field perturbation $B_0$ versus time at different axial positions $z$ from the electrode at which a current pulse $I_e(t) = 2 \text{ A}$ is drawn for $\Delta t = 28 \mu\text{sec}$. Both turn-on and turn-off of the current produces a propagating oscillatory magnetic field perturbation. The wave packet is characterized by a falling "tone" characteristic for low-frequency whistlers whose group and phase velocities vary with frequency as $v_g = 2 v_{ph} = 2c (\omega_B e^{-1/2})/v_p$.

![Magnetic field $B_0(t)$ vs. time at different axial positions $z$ from the disc electrode ($r = 2.5 \text{ cm}$) collecting an electron current $I_e(t)$ shown in the bottom trace. Both the turn-on and turn-off of the current pulse excite a propagating wave burst of falling tone characteristic of low-frequency whistlers.](image)

In order to establish a well-defined frequency $\omega$ rather than a spectrum of $\omega$ characteristic for a single pulse, the current $I_e(t)$ was square-wave modulated at a frequency $\omega$ and the corresponding oscillatory magnetic field components measured as shown in Figure 7. Figure 7a shows the axial variation of $B_0(z)$ which exhibits a well defined wavelength ($\lambda_z = 20 \text{ cm}$) and a phase velocity $v_{ph} = 3.7 \times 10^5 \text{ cm/sec}$ with direction of wave propagation along $B_0$ away from the electrode. Two orthogonal magnetic field components $B_x$, $B_y$ are $90^\circ$ out of phase such that the vector magnetic field $B = (B_x, B_y)$ forms in space a right-hand circular spiral as shown in Figure 7b. When the frequency $\omega$ is varied and the axial wavenumber $k_z = 2\pi/\lambda_z$ is plotted the
Fig. 7. Wave analysis for repetitive current pulses (top trace). (a) Interferometer traces \( B_x(z) \) at different delay times \( \tau \) in the reference signal, \( \cos(\omega(t-\tau)) \). The traces indicate undamped waves propagating away from the electron-collecting electrode at \( z = 0 \) with well-defined phase-velocity \( v_b \). (b) Interferometer traces of two orthogonal field components \( (B_x, B_y) \) vs. \( z \) which show that the field vector \( \mathbf{B} \) describes a right-hand circular spiral (see insert and corresponding points A, B, C, D). (c) Dispersion diagram \( \omega \) vs. \( k_1 \), suitably normalized, showing that the data points fall within measurement accuracy on the theoretical dispersion curve for whistler waves, \( \left(\frac{k_1 c}{\omega}\right)^2 = 1 - \omega_p^2 / \omega (\omega c_e) \) with \( \omega_p = 2 \pi \times 8 \text{ GHz}, \omega c_e = 2\pi \times 20 \text{ MHz}. \)

Data points fall on to the theoretical dispersion curve \( \omega(k_1) \) for whistlers normalized to the experimental parameters, shown in Fig. 7c. Thus, from the measurement of phase velocity, polarization and dispersion it is shown conclusively that oscillating currents excite whistler waves, and that switched currents penetrate along \( B_0 \) with the speed of a whistler wave.
packet. Thus, in the regime of unmagnetized ions the electrons and associated waves determine transport and dissipation of magnetic energy.

**MAGNETIC NEUTRAL SHEETS**

In order to produce in a laboratory plasma the classical two-dimensional magnetic field topology for reconnection /24/ the setup of Figure 8 has been used /25/. A linear plasma column \( n_e = 10^{12} \) cm\(^{-3} \), \( T_e = 10 \) keV, Ar, He) is produced by a pulsed dc discharge (50 V, 1000 A, 5 msec) between a 1 m diam. oxide coated cathode and an end anode, confined by a weak axial magnetic field \( B_Y = 20 \) G. Two parallel-plate electrodes are placed adjacent to but insulated from the plasma column, carrying pulsed axial currents \( I_1(t) \) increasing in time \((dt = 100 \) µsec, \( I_1 = 10 \) kA). All currents are closed in a coaxial fashion via the conducting outer cylindrical chamber wall. Without the plasma the magnetic field of the plate current system has an X-type magnetic null point on axis, as shown in the cross-sectional view on the right hand side of Figure 8. However, in the presence of plasma between the plates a large axial plasma current \( I_a = 1000 \) A is induced whose 0-type field topology superimposes on the X-type field of the plate currents and can result in the formation of a neutral sheet \( B_z = 0 \), \( B_x, B_y \neq 0 \). The induced plasma current is mainly carried by electrons streaming axially from cathode to anode and, to a small degree, by ions streaming into the cathode \( (i_e/m_e)/\gamma_m = 0.3\% \).

**EXPERIMENTAL ARRANGEMENT**

![Figure 8. Schematic diagram of the experimental setup to generate a magnetic neutral sheet in a discharge plasma column (left: side view; right: cross-sectional view).](image)

Figure 9a shows a measurement of the vector magnetic field \( (B_x, B_z) \) vs. \( x, z \) for a typical neutral sheet configuration. In order to facilitate comparison between the laboratory and magnetotail geometries the conventional coordinate system has been adopted where \( x \) is along the neutral sheet, \( z \) is normal to the sheet, and \( y \) is the coordinate along the neutral line or current flow. Note that in the laboratory set up there is a magnetic field component \( B_Y = 20 \) G comparable to the transverse fields \((B_x, B_z) \leq 15 \) G). Characteristics of the observed neutral sheet are shown in Figure 9b which displays \( B_x(x=0), \) the current density profile \( j_y(x=0) = (3B_y/2b)/\mu \) and various theoretical scale lengths. Once can see that the half width of the current sheet, \( \Delta_y = 5 \) cm, is small compared to characteristic ion scale lengths \((c/\omega_i) = 60 \) cm, \( r_{ci} = v_i/\omega_i = 12 \) cm in He), but larger than the corresponding electron scale lengths \((c/\omega_p) = 0.7 \) cm, \( r_{ce} = 0.4 \) cm) and the size of the theoretical diffusion region \((\delta_{diff} = 0.3 \) cm for \( R_m = 1 \)). Since MHD theory assumes scale lengths large compared to an ion Larmor radius it cannot account for the physics near a neutral sheet at distances \( z \leq r_{ci} \). For example, magnetic field and fluid are not frozen-in up to the diffusion region where reconnection is allowed. As the ions become demagnetized \((B/\Omega_B = r_{ci})\) the magnetic field is decoupled from the ion inertia and free to follow the dynamic electrons. Thus, magnetic energy can be transported at the fast whistler rather than the slow Alfvén velocity \((v_w/v_A = (m_i/m_e)^{1/2}) \). In the present case of a significant field component \( B_Y \) the energy flow is along the separator rather than in the x-z plane. When the current sheet is perturbed the magnetic perturbations travel along y as fast whistler modes rather than in the x-z plane as slow Alfvén waves /25/. There are, of course, no Petschek shocks in this regime /26/. The fast energy transport in the y-direction emphasizes the importance of considering a realistic three-dimensional geometry which does not assume infinite length in the y-direction. We will come back to this aspect when considering the process of conversion of magnetic energy into particle kinetic energy.
Fig. 9. (a) Vector magnetic field \(B_x, B_z\) in the x-z plane at \(y = 137\) cm from the cathode showing a classical neutral sheet topology, however with \(B_x = 20\) G normal to the plane. (b) Magnetic field \(B_x\) vs. \(z\) at \(y = 0\) and its derivative \(J_y(z)\). Scaling parameters to the left show that the current sheet thickness \(\Delta_y = 5\) cm is smaller than the ion Larmor radius but larger than characteristic electron scale lengths. MHD concepts such as frozen-in fields, Alfvén waves and magnetic Reynolds numbers are not meaningful near such neutral sheets.

It was pointed out that the ions play a negligible role compared to electrons in contributing directly to currents and magnetic fields. However, in a plasma of unmagnetized ions and magnetized electrons space-charge electric fields arise which are as important to consider as magnetic and pressure forces. Figure 10 shows measurements of the ion flow velocity field.

Fig. 10. Flow velocity field of the ions normalized to the sound speed, \((v_x, v_z)/c_s\), in the x-z plane showing that at late times, \(t = 80\) µsec, the unmagnetized ions are accelerated by space-charge electric fields to the same motion as the current-carrying magnetized electrons subject to the magnetic \(J \times B\) force.
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which indicate that unmagnetized ions are accelerated in the x-z plane to the classical MHD flow pattern across the separatrix /24/. This acceleration must arise from space-charge electric fields since the magnetic J x B force is mainly exerted on the current-carrying magnetized electrons. The ion jetting velocity is not related to or limited by the Alfvén speed as predicted by MHD theories. Electron E x B drifts cause Hall currents /28/ which are not considered in MHD models. Thus, on the scale length of an ion Larmor radius neutral sheets show a rather different behavior than extrapolated from MHD theories or simulations.

TEARING INSTABILITIES

The current sheet of Figure 9 with approximate scales \( \Delta x = 45 \text{ cm}, \Delta z = 5 \text{ cm} \) is observed to be stable to tearing as long as the external plate current is increased in time (\( \Delta t < 100 \mu \text{sec} \)). Since the tearing instability strongly depends on the ratio of current sheet thickness to length \( \Delta x/\Delta z \), we have reduced this parameter and observed the onset of tearing. Figure 11 shows the experimental arrangement /22/. Between cathode and anode a grid with rectangular

\[ \text{VARIATION OF CURRENT SHEET THICKNESS} \]

\[ \text{ANODE} \quad \text{CONSTRING GRID} \quad \text{CATHODE} \]

\[ \begin{aligned}
B_x(t) & \quad B_{y0} \\
I_{\text{pl.}} & \quad E_y \\
I_{\text{ext.}} & \quad 200 \text{ V}
\end{aligned} \]

\[ \begin{aligned}
\text{CONSTRINGED} & \quad \text{NORMAL}
\end{aligned} \]

Fig. 11. Experimental setup to vary the thickness of the current sheet in order to study tearing instabilities. The electron current is forced to flow through a narrow (\( \Delta z = 1 \text{ cm} \)), long (\( \Delta x = 75 \text{ cm} \)) aperture in a negatively biased grid located between cathode and anode (\( y = 100 \text{ cm} \)). Measured current density profiles \( J(z) \) with and without constriction are shown in the bottom two graphs. Spontaneous tearing is observed for the narrow current sheet (\( \Delta z = 1 \text{ cm} \)).
iris ($\Delta_x = 1$ cm, $\Delta_y = 75$ cm) is inserted and pulsed to a large negative potential so as to repel all electrons except those streaming through the iris. The current sheet is thereby compressed from $\Delta_x = 3$ cm to $\Delta_x = 1$ cm as indicated by the current density profiles $J_y(z)$ in Figure 11. The spontaneous growth of magnetic islands has been observed for such thin current sheets /25/. The growth rate of the instability is found to be consistent with that of collisionless electron tearing modes /29/. Further magnetic field measurements in several x-z planes at different axial distances $\Delta y$ from the iris are displayed in Figure 12. It is found

Fig. 12. Magnetic field lines projected in x-z planes at different distances $\Delta y$ from the injection point of a narrow current sheet ($\Delta_x = 1$ cm). The current sheet is subject to a fast electron tearing mode producing multiple X and O points. Note that tearing and merging of islands occur periodically along y with wavenumber $k_y = k_x = 2m/(120$ cm), i.e. the tearing mode is three-dimensional.
that tearing and merging occur periodically along \( y \) with wavenumber \( k_y = k_y = 2\pi/20 \text{ cm}^{-1} \). The axial propagation speed of magnetic perturbations is that of whistler waves. These observations demonstrate clearly that magnetic instabilities and topologies in neutral sheets with \( B_z \neq 0 \) are controlled by the fast speed of electrons and whistlers and are, in general, three-dimensional.

ENERGY CONVERSION

One of the most important issues in magnetic reconnection is the mechanism of conversion of magnetic field energy into particle kinetic energy. Classical MHD models predict that magnetic energy is mainly converted into streaming ion (fluid) energy via \( \mathbf{J} \times \mathbf{B} \) forces at current sheets or shocks /31/. However, for current sheets which are not many ion Larmor radii thick the MHD model is questionable. The analysis of the present laboratory experiment shows that the dominant energy transfer is entirely different than envisioned by standard models.

Figure 13 presents a simplified diagram of the global current system. An external energy source drives a time-varying current \( I_1(t) \) through a primary circuit. By Faraday's law the time rate of change of the reconnected linked magnetic flux \( \Phi \) represents an induced voltage \( V = -d\Phi/dt \) in the secondary circuit containing the discharge plasma. A secondary current \( I_2(t) \) is induced. Dissipation arises due to irreversible particle acceleration by the induced electric field in the secondary circuit. Only its line integral \( V = \int \mathbf{E} \cdot d\mathbf{l} \) is known, the local electric field \( \mathbf{E} \) depends on the local properties of the secondary circuit. In the laboratory experiment the return current path through the 5 cm thick aluminum chamber wall presents a negligible resistance, hence voltage drop \( (dV/V < 1\%) \). The plasma column itself is a good conductor along which only 20\% of the induced voltage \( (V = 100 \text{ Volt}) \) drops off.

The bulk of the induced voltage in this laboratory device drops off at the cathode sheath, a thin (\( d < 1 \text{ mm} \)) region of strong electric fields \( (E > 1000 \text{ V/cm}) \) in which emitted electrons are accelerated into the plasma and ions into the cathode. Although the sheath is entirely collisionless it produces effective dissipation of the induced electromagnetic energy in the circuit and presents a classical case of inertial reconnection /32/. The sheath is an
inherent, unavoidable part of this and similar laboratory setups; without it there would be no electron emission, the current would drop by two orders of magnitude to the ion saturation current, the magnetic field topology and reconnection rate would be essentially the same as in vacuum.

The accelerated particles from the sheath have been measured directly with a directional, Faraday-type energy analyzer /33/. Figure 14a displays the electron current vs. particle energy for electrons streaming away from (top trace) and toward (bottom trace) the cathode. There is a large flux of energetic (up to 120 eV) electrons emitted by the cathode which cannot be accounted for by the dc discharge ($V_{dis} = 40$ V) but arises from the added reconnection voltage at the cathode sheath. By measuring the flux at all angles the electron distribution function in three-dimensional velocity space can be obtained and displayed as surfaces of constant $f(v_x, v_y, v_z)$, an example of which is shown in Figure 14b. The tail electrons are clearly visible as the deviation from the spherical surface associated with the Maxwellian background electrons. The energy imparted by reconnection to tail electrons is partly transferred to the bulk electrons by anomalous heating through beam-plasma instabilities. Most of the converted electromagnetic energy is transferred to particle kinetic energy, which is dissipated at the axial boundaries, relatively little is transferred to transverse $(x,z)$ fluid motion by jetting.

Fig. 14. Electron energization by inertial reconnection producing anisotropic distributions. (a) Electron flux to directional Faraday-cup vs. energy. From the cathode a tail of 120 eV electrons is injected into the plasma. The energetic electrons are accelerated at the cathode sheath where most of the reconnection voltage drops off rather than in the bulk of the plasma where the electrons are essentially Maxwellian. (b) Electron distribution function in three-dimensional velocity space displayed as a contour $f(v_x, v_y, v_z) = \text{const}$. The energetic electrons form the protrusion from the spherical surface of the Maxwellian bulk electrons. The free energy of the tail excites various microinstabilities in the neutral sheet.

Other situation of localized electric fields along the separator have been observed. At high current densities (electron $v_d/v_e = 1$) a nonstationary double layer with strong $E_y \parallel B_0$ is formed spontaneously /34/. A strong normal magnetic field $B_z \gg B_0$ impedes the electron current $J_y$ such that the reconnection voltage drops off across the region of strong $B_z$ /25/.
Any constrictions, boundary layers or nonuniformities in the current circuit are potential regions of localized electric fields and dissipation. The two-dimensional picture of reconnection where the induced electric field $E_y = \beta_\phi / \partial t$ is constant along $y$ is highly idealized and may be applicable only to highly symmetric cases such as toroidal current systems (tokamaks).

CONCLUSIONS

The present laboratory experiment has focussed on reconnection processes in the regime of non-magnetized ions where MHD theories break down and little theoretical guidance is available. A magnetic field component $B_y$ along the separator and three-dimensional effects ($\beta_\phi \neq 0$) are present. From the various observations it is concluded that magnetic fields are highly dynamic compared to MHD predictions. Magnetic energy travels in the form of fast whistler waves along the separator rather than as slow Alfvén waves in the $x$-$z$ plane. Nonuniformities in the global current circuit lead to localized reconnection electric fields $E_y$ and collisionless particle energization (inertial reconnection). Since energy storage and dissipation may be spatially separated along $y$ transport processes along the separator are very important.

How are the results from these laboratory experiments related to the present picture of reconnection in magnetospheric and solar physics? The basic assumption in most reconnection models is the validity of ideal MHD with the exception of the null region which is approximated as a non-ideal MHD region (diffusion region). Measurements of the current sheet thickness of the magnetotail during plasma sheet thinning indicate that current sheets are as narrow as an ion gyroradius do exist. Detailed measurements within these sheets are, however, complicated by random tail motions. Current sheets in solar plasmas cannot be resolved with present observational techniques but as instruments improve more fine-scale structures have usually been observed. Since the photospheric plasma and this laboratory plasma have similar parameters it is not impossible that the fine-scale structure in electron current sheets is the same in both cases. In those situations where the magnetic field is decoupled from the ions the dynamics is fast, electron energization dominates, and transport along a neutral line with $B_y$ component is very important. New theoretical models should be developed to account for the physics of reconnection in thin current sheets.

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